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Seesaw-Type Quark and Lepton Mass Matrices and U(3)-Family Nonet Higgs Bosons

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Abstract

A unified mass matrix model of quarks and leptons with a seesaw-type form $M_f = m M_F^{-1} m$ is proposed on the basis of a non-standard Higgs scenario. The matrix m is provided by U(3)-family nonet bosons ϕ, and the matrix M_F is a mass matrix of heavy fermions F_i corresponding to the ordinary fermions $f_i = \nu_i, e_i, u_i, d_i$ $(i = 1, 2, 3)$. It is shown that a Higgs potential of ϕ with a broken U(3)-family symmetry leads to a desirable charged lepton mass formula when $M_E \propto \mathbf{1}$. Then, phenomenologically desirable forms of heavy quark mass matrices M_Q $(Q = U, D)$ are investigated.

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Recent measurements of tau lepton mass m_τ have found [1] that the observed value excellently satisfies a charged lepton mass formula

$$m_e + m_\mu + m_\tau = \frac{2}{3}(\sqrt{m_e} + \sqrt{m_\mu} + \sqrt{m_\tau})^2 \ , \tag{1}$$

which provides $m_\tau = 1777$ MeV for the input values of m_e and m_μ. The relation (1) has first been speculated on the basis of a composite model [2] and then an extended technicolor-like model [3] by one of the authors (Y.K.). In those models, it is essential to assume that the mass matrix M_f for fermions f_i ($i = 1, 2, 3$: family indices) is given by the form

$$M_f = m_0^f \, G O_f G \ , \tag{2}$$

where $G = \mathrm{diag}(g_1, g_2, g_3)$. If we assume that the parameters g_i satisfy the relations $g_i = g_i^{(8)} + g^{(1)}$, $\sum_i g_i^{(8)} = 0$ and $\sum_i \left(g_i^{(8)}\right)^2 = 3\left(g^{(1)}\right)^2$, and the matrix form of O_f is given by $O_e = \mathbf{1}$ ($\mathbf{1}$ is a 3×3 unit matrix) for the charged lepton mass matrix M_e, then we can obtain the relation (1). Furthermore, one of the authors (Y.K.) has recently pointed out [4] that the mass matrix (2) has a possibility that it can provide a unified description of quark and lepton mass matrices by considering a special form of O_f (without re-fitting the parameters g_i for each $f = u, d, \nu$).

In Ref.[3], the matrix form (2) was speculated on the basis of an extended technicolor-like model. Then, in that model, there were too many fermions, because we must consider many technicolored fermions $F_{i\alpha}$ (α is technicolor index) corresponding to the ordinary fermions f_i. Therefore, in that model, SU(3)-color cannot be asymptotic free.

In the present paper, we propose an alternative model based on a non-standard Higgs model, which leads to a seesaw-type mass matrix $M_f \simeq m M_F^{-1} m$ and in which U(3)-family nonet Higgs scalars ϕ play an essential role of deriving the mass formula (1). First, in order to give an outline of the model, we will discuss the case for charged leptons, where M_F ($F = E$) is simply assumed as $M_E \propto \mathbf{1}$. Next, we will discuss possible forms of the heavy fermion mass matrices M_F ($F = U, D, N$) from the phenomenological point of view (but within the context of our Higgs scenario).

In our scenario, we prepare the following fermions: $f = \ell, q$ ($\ell = (\nu, e)$, $q = (u, d)$) and $F = N, E, U, D$, which belong to $f_L = (\mathbf{2}, \mathbf{1}, \mathbf{3})$, $f_R = (\mathbf{1}, \mathbf{2}, \mathbf{3})$, $F_L = (\mathbf{1}, \mathbf{1}, \mathbf{3})$, and $F_R = (\mathbf{1}, \mathbf{1}, \mathbf{3})$ of SU(2)$_L \times$SU(2)$_R \times$U(3)$_{family}$, respectively.

Up- and down-heavy fermions, F^{up} and F^{down}, are distinguished by hypercharge Y (note that $Y \neq B-L$ for the heavy fermions): Hypercharges of the heavy fermions (N, E) and (U, D) take the values $(0, -2)$ and $(4/3, -2/3)$, respectively. In the present model, differently from the standard $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_{B-L}$ model, we do not consider Higgs scalar fields which belong to $(\mathbf{2}, \mathbf{2})$ of $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$, so that there are no Higgs fields which couple with $\overline{f}f$ at tree level. We assume only the following Yukawa interactions:

$$H_{Yukawa} = g_F \sum_{i,j} \overline{F}^i (\Phi_F)_i^j F_j + g_L \sum_{i,j} \left(\overline{f}^i_L (\phi_L)_i^j F^{down}_{Rj} + \overline{f}^i_L (\widetilde{\phi}_L)_i^j F^{up}_{Rj} + H.C. \right) + (L \leftrightarrow R) \ , \tag{3}$$

where $\phi = (\phi^+, \phi^0)$ and $\widetilde{\phi} = (\overline{\phi}^0, -\phi^-)$. The scalar fields ϕ_L and ϕ_R belong to $(\mathbf{2}, \mathbf{1}, \mathbf{8}+\mathbf{1})$ and $(\mathbf{1}, \mathbf{2}, \mathbf{8}+\mathbf{1})$ of $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(3)_{family}$, respectively, and the vacuum expectation values (VEV's) $\langle \phi^0_L \rangle$ and $\langle \phi^0_R \rangle$ provide left- and right-handed weak boson masses $m(W_L)$ and $m(W_R)$, respectively. The fields Φ_F which belong to $(\mathbf{1},\mathbf{1},\mathbf{8}+\mathbf{1})$ do not contribute to weak boson masses $m(W_L)$ and $m(W_R)$, but play only a role of providing extremely large masses for vector-like fermions F. (The structure of $\langle \Phi_F \rangle$ will be discussed later.) Under the approximation of $M_F \gg m_L, m_R$ ($M_F = g_F \langle \Phi_F \rangle$, $m_L = g_L \langle \phi^0_L \rangle$, and so on), we obtain a seesaw-type mass matrix form $M_f \simeq m_L M_F^{-1} m_R$.

First, we discuss the charged lepton mass matrix M_e. We assume that charged heavy leptons E couple only with a $\mathrm{U}(3)_{family}$ singlet scalar field Φ_0, which belongs to $(\mathbf{1}, \mathbf{1}, \mathbf{1})$ of $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times$ $\mathrm{U}(3)_{family}$, so that the matrix form O_e in (2) is given by $O_e = \mathbf{1}$. On the other hand, for the scalar fields ϕ_L, we assume the following Higgs potential which is approximately invariant under the U(3)-family symmetry:

$$V(\phi) = \mu^2 \mathrm{Tr}\left(\phi^- \phi^+ + \overline{\phi}^0 \phi^0\right) + \frac{1}{2}\lambda \left[\mathrm{Tr}\left(\phi^- \phi^+ + \overline{\phi}^0 \phi^0\right)\right]^2$$

$$+\frac{1}{2}\lambda' \left[\mathrm{Tr}\left(\phi^- \phi^-\right) \mathrm{Tr}\left(\phi^+ \phi^+\right) + 2\mathrm{Tr}\left(\phi^- \overline{\phi}^0\right) \mathrm{Tr}\left(\phi^+ \phi^0\right) + \mathrm{Tr}\left(\overline{\phi}^0 \overline{\phi}^0\right) \mathrm{Tr}\left(\phi^0 \phi^0\right)\right]$$

$$+\eta \left(\phi^-_1 \phi^+_1 + \overline{\phi}^0_1 \phi^0_1\right) \mathrm{Tr}\left(\phi^-_8 \phi^+_8 + \overline{\phi}^0_8 \phi^0_8\right) + \eta' \mathrm{Tr}\left[\left(\phi^-_1 \phi^+_8 + \overline{\phi}^0_1 \phi^0_8\right)\left(\phi^-_8 \phi^+_1 + \overline{\phi}^0_8 \phi^0_1\right)\right] \ , \tag{4}$$

where, for convenience, we have omitted the index L. Here, the traces are taken over the family indices and $\phi_8 = (\phi^+_8, \phi^0_8)$ and $\phi_1 = (\phi^+_1, \phi^0_1)$ denote octet and

singlet components of ϕ, respectively: $\phi = \phi_8 + (1/\sqrt{3})\phi_1 \mathbf{1}$. In the potential (4), the U(3)-family invariance of $V(\phi)$ is explicitly broken by the η- and η'-terms, while an SU(3)-family symmetry is still unbroken. Hereafter, we sometime use the language of SU(3)-family instead of broken U(3)-family. Of course, the potential (4) is not a general form of the $\mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{SU}(3)_{family}$ invariant potential.

For $\mu^2 < 0$, conditions for minimizing the potential $V(\phi)$ are as follows:

$$\left[\mu^2 - \lambda \mathrm{Tr}(\mathbf{v}^\dagger \mathbf{v}) - (\eta + \eta')\mathrm{Tr}(\mathbf{v}_8^\dagger \mathbf{v}_8)\right] v_1^* - \lambda' \mathrm{Tr}(\mathbf{v}^\dagger \mathbf{v}^\dagger) v_1 = 0 \ , \tag{5a}$$

$$\left[\mu^2 - \lambda \mathrm{Tr}(\mathbf{v}^\dagger \mathbf{v}) - (\eta + \eta') v_1^* v_1\right] (\mathbf{v}_8^\dagger)_i^j - \lambda' \mathrm{Tr}(\mathbf{v}^\dagger \mathbf{v}^\dagger)(\mathbf{v}_8)_i^j = 0 \ , \tag{5b}$$

and equations exchanged as ($\mathbf{v} \leftrightarrow \mathbf{v}^\dagger$, $\mathbf{v}_8 \leftrightarrow \mathbf{v}_8^\dagger$, and $v_1 \leftrightarrow v_1^*$) in (5$a$) and (5$b$), respectively, where $\mathbf{v} = \langle \phi^0 \rangle$, $\mathbf{v}_8 = \langle \phi_8^0 \rangle$ and $v_1 = \langle \phi_1^0 \rangle = (\mathrm{Tr}\,\mathbf{v})/\sqrt{3}$. These conditions lead to the relations $\mathbf{v}^\dagger = \mathbf{v}$ and

$$v_1^* v_1 = \mathrm{Tr}\left(\mathbf{v}_8^\dagger \mathbf{v}_8\right) = -\mu^2 / \left[2(\lambda + \lambda') + \eta + \eta'\right] \ ,$$

so that we obtain the relation

$$\mathrm{Tr}\left(\mathbf{v}^2\right) = \frac{2}{3} \left(\mathrm{Tr}\,\mathbf{v}\right)^2 \ . \tag{6}$$

We assume that a Higgs potential $V(\phi_R)$ has the same structure with $V(\phi_L)$, i.e., each term in $V(\phi_R)$ takes the coefficient which is exactly proportional to the corresponding term in $V(\phi_L)$. This assumption means that there is a kind of "conspiracy" between $V(\phi_R)$ and $V(\phi_L)$. However, in this paper, we do not go into this problem moreover. When we assume $\langle \phi_R^0 \rangle \propto \langle \phi_L^0 \rangle$, we can obtain the mass formula (1) from the relation (6).

Re-derivation of the relation (1) on the basis of a Higgs potential with a broken U(3)-family symmetry has also been done by one of the authors (Y.K.) [5]. However, since his potential lacked the λ'- and η'-terms in (4), too many massless Nambu-Goldstone states appeared, although he successfully derived the relation (6). In order to avoid such excess of massless states, we need the λ'-term. On other hand, if we add, for example, a term $\mathrm{Tr}\left[(\phi^- \phi^+ + \overline{\phi}^0 \phi^0)^2\right]$ to the potential (4), we cannot obtain the relation (6). In order to the relation (6), the potential $V(\phi)$ is restricted to a special form under the broken U(3)-family symmetry.

Next, we discuss the form of O_f. In Ref.[4], it was pointed out that a form $O_f = \mathbf{1} + 3a_f X(\phi_f)$ can provide successful predictions of quark masses and

Kobayashi-Maskawa (KM) [6] matrix elements, where the matrix $X(\phi_f)$ is a democratic type matrix with a phase factor ϕ_f,

$$X(\phi) = \frac{1}{3} \left(\begin{array}{ccc} 1 & e^{i\phi} & 1 \\ e^{-i\phi} & 1 & 1 \\ 1 & 1 & 1 \end{array} \right) . \tag{7}$$

However, in the context of a Higgs scenario, it is not so easy to derive the matrix form of the second term $X(\phi_f)$, which is not a rank one matrix for $\phi_f \neq 0$. In the present paper, we consider alternative form of O_f.

We assume that the U$(3)_{family}$ singlet field Φ_0 couples with all heavy fermions F universally. In addition to Φ_0, we assume a Higgs field Φ_X which couples only with a heavy quark states $F_S = (F_1 + F_2 + F_3)/\sqrt{3}$ ($F_i = U_i, D_i$) (a symmetric representation of the permutation group S_3 [7]). The VEV $\langle \Phi_X \rangle$ provides a democratic mass matrix term without phase factors, $X(0)$, which is a rank one matrix. We consider that the coupling constants of Φ_X with heavy fermions F are different according as $F = U$ or $F = D$. Therefore, the matrix form O_f in (2) is given by

$$O_f = \mathbf{1} + 3 a_f e^{i\alpha_f} X(0) , \tag{8}$$

i.e.,

$$M_F \propto O_f^{-1} = \mathbf{1} + 3 b_f e^{i\beta_f} X(0) , \tag{9}$$

where

$$b_f e^{i\beta_f} = -a_f e^{i\alpha_f}/(1 + 3 a_f e^{i\alpha_f}) . \tag{10}$$

The reason that we still consider a democratic matrix from in O_f is motivated by only a phenomenological reason suggested in Ref.[4], i.e., by the fact that for up-quark mass matrix with $\beta_u = 0$, we can obtain the successful mass relation [4]

$$m_u/m_c \simeq 3 m_e/4 m_\mu , \tag{11}$$

for a small value of $\varepsilon_u \equiv 1/a_u$. Note that the ratio m_u/m_c is insensitive to the parameter a_u. The parameter $\varepsilon_u \equiv 1/a_u$ is determined by the mass relation

$$m_c/m_t \simeq 2(m_\mu/m_\tau)\varepsilon_u .$$

Differently from the model given in Ref.[4], down-quark mass matrix M_d with $\alpha_d \neq 0$ in the present model is not Hermitian. We will demonstrate that the

present model with the form (8) also can provide reasonable predictions of quark mass ratios and KM matrix by adjusting our parameters a_u, a_d and α_d.

In the present model, a case $a_d \simeq -0.5$ can provide phenomenologically interesting predictions as seen below. For small values of $|\alpha_d|$ and $\varepsilon_d \equiv -(2+a_d^{-1})$, we obtain the down-quark mass ratios

$$m_s/m_b \simeq (1/2)\kappa \left(1 - 48\sqrt{2\varepsilon/3}\right) \ , \tag{12}$$

$$m_d/m_s \simeq 16(\varepsilon/\kappa^2) \left(1 + 96\sqrt{2\varepsilon/3}\right) \ , \tag{13}$$

where

$$\kappa = \sqrt{\sin^2 \frac{\alpha_d}{2} + \left(\frac{\varepsilon_d}{4}\right)^2} \ , \quad \varepsilon = \frac{m_e m_\mu}{m_\tau^2} \ . \tag{14}$$

We also obtain

$$m_d m_s / m_b^2 \simeq 4 m_e m_\mu / m_\tau^2 \ , \tag{15}$$

as a relation which is insensitive to the small parameters $|\alpha_d|$ and ε_d.

Furthermore, we can obtain ratios of up-quarks to down-quarks, for example,

$$m_u/m_d \simeq 6\kappa \sim 12 m_s/m_b \ . \tag{16}$$

Suitable choice of small values of ε_d and α_d ensures $m_u/m_d \sim O(1)$ in spite of $m_t \gg m_b$. From (10), a small value $|\varepsilon_u| = 1/|a_u| \simeq 0$ means $b_u \simeq -1/3$, while a small value $|\varepsilon_d| = |2 + a_d^{-1}| \simeq 0$ means $b_d \simeq -1$. It is noted that, in spite of the large ratio of m_t/m_b, the ratio of b_d/b_u is not so large, i.e., $b_d/b_u \simeq 3$.

Then, let us discuss the KM matrix elements V_{ij}. The KM matrix V is given by $V = U_L^u P U_L^{d\dagger}$, where U_L^u and U_L^d are defined by $U_L^u M_u M_u^\dagger U_L^{u\dagger} = \mathrm{diag}(m_u^2, m_c^2, m_t^2)$ and $U_L^d M_d M_d^\dagger U_L^{d\dagger} = \mathrm{diag}(m_d^2, m_s^2, m_b^2)$, respectively, and P is a phase matrix. Here, we have considered that the quark basis for the mass matrix (2) can, in general, deviate from the quark basis of weak interactions by some phase rotations, The simplest case $P =$diag$(1,1,1)$ cannot provide reasonable predictions of $|V_{ij}|$. Only when we take $P = \mathrm{diag}(1,1,-1)$, we can obtain reasonable predictions for both quark mass ratios and KM matrix elements, although it is an open question why such a phase inversion is caused on the third family quark. The predictions of $|V_{ij}|$ are sensitive to every values of ε_u, ε_d and α_d, so that it is not adequate to express $|V_{ij}|$ as simple approximate relations such as those in (11)–(16). Therefore, we

will show only numerical results for $|V_{ij}|$. For example, by taking $b_u = -0.3295$, $b_d = -1.072$, $\beta_u = 0$ and $\beta_d = 18.5^\circ$, which are chosen by fitting the quark mass ratios, we obtain the following predictions of quark masses, KM matrix elements $|V_{ij}|$ and the rephasing-invariant quantity [8] J:

$$\begin{array}{lll} m_u = 0.00228 \text{ GeV} \ , & m_c = 0.591 \text{ GeV} \ , & m_t = 170 \text{ GeV} \ , \\ m_d = 0.00429 \text{ GeV} \ , & m_s = 0.0875 \text{ GeV} \ , & m_b = 3.02 \text{ GeV} \ , \end{array} \tag{17}$$

$$\begin{array}{c} |V_{us}| = 0.223 \ , \quad |V_{cb}| = 0.0542 \ , \quad |V_{ub}| = 0.00309 \ , \quad |V_{td}| = 0.0146 \ , \\ |V_{ub}/V_{cb}| = 0.0570 \ , \quad J = 2.30 \times 10^{-5} \ . \end{array} \tag{18}$$

The prediction $|V_{cb}| = 0.0542$ in (18) is somewhat large in comparison with the experimental value $|V_{cb}| = 0.043 \pm 0.007$ [9]. If we use $P = (1, 1, -e^{i\delta})$ with a small phase value δ instead of $P = (1, 1, -1)$, we can obtain more excellent predictions without changing predictions of quark masses in (17): for example, when we take $\delta = -3.4^\circ$, we obtain $|V_{us}| = 0.223$, $|V_{cb}| = 0.0431$, $|V_{ub}| = 0.00282$ ($|V_{ub}/V_{cb}| = 0.0654$), $|V_{td}| = 0.01184$ and $J = 1.69 \times 10^{-5}$.

In the numerical predictions of quark masses, (17), we have used a common enhancement factor of quark masses to lepton masses, $m_0^u/m_0^e = m_0^d/m_0^e = 3$, in order to compare with quark mass values [10] at the energy scale $\mu = \Lambda_W \equiv (\mathrm{Tr}\langle\phi_L^0\rangle^2)^{1/2}/\sqrt{2} = (\sqrt{2}G_F)^{-1/2}/\sqrt{2} = 174$ GeV: $m_u = 0.0024 \pm 0.0005$ GeV, $m_d = 0.0042 \pm 0.0005$ GeV, $m_c = 0.605 \pm 0.009$ GeV, $m_s = 0.0851 \pm 0.014$ GeV, $m_t = 174 \pm 16$ GeV, and $m_b = 2.87 \pm 0.03$ GeV, where we have used [9] $\Lambda_{\overline{MS}}^{(4)} = 0.26$ GeV. Although we are happy if we can explain such the factor $m_0^q/m_0^e = 3$ by evolving quark and lepton masses from $\mu = \Lambda_X$ (a unification scale [11]) to $\mu = \Lambda_W$, unfortunately, it is not likely to derive such a large factor ~ 3 from the conventional renormalization calculation. We must assume an additional enhancement mechanism, for example, a different coupling strength of Φ_0 with heavy quarks from that with heavy leptons.

For neutrinos, if we take $b_\nu \simeq -1/3$ and $\beta_\nu = 0$ similar to the case of up-quarks, we can obtain an interesting prediction [12] for the neutrino mixing between ν_e and ν_μ,

$$\sin\theta_{e\mu} \simeq (1/2)\sqrt{m_e/m_\mu} \simeq 0.035 \ ,$$

which is in good agreement with the value $\sin\theta_{e\mu} \simeq 0.04$ ($\sin^2 2\theta_{e\mu} \simeq 7 \times 10^{-3}$) suggested by the GALLEX data [13]. However, in the present stage, we do not

have any unified understanding of b_f and β_f, i.e., they are nothing more than phenomenological parameters.

Finally, we comment on physical Higgs bosons ϕ_L in the present scenario. We define three mixing states among ϕ_1^1, ϕ_2^2 and ϕ_3^3 as follows:

$$\phi_3' = (1/\overline{v})\left(v_1\phi_1^1 + v_2\phi_2^2 + v_3\phi_3^3\right) \ ,$$

$$\phi_1' = (1/\overline{v})\sqrt{2/3}\left[(v_3 - v_2)\phi_1^1 + (v_1 - v_3)\phi_2^2 + (v_2 - v_1)\phi_3^3\right] \ ,$$

and ϕ_2' are a state which is orthogonal to ϕ_1' and ϕ_3', where $\langle\phi^0\rangle = \mathrm{diag}(v_1, v_2, v_3)$ and $\overline{v} = (v_1^2 + v_2^2 + v_3^2)^{1/2}$. Moreover, for convenience, we rewrite ϕ^+ and ϕ^0 as

$$\phi^+ = i\Pi^+ \ , \quad \phi^0 = (H^0 - i\Pi^0)/\sqrt{2} \ ,$$

Then, for the $\Pi^\pm$ and Π^0 states, the spontaneous symmetry breaking $\langle\phi^0\rangle \neq 0$ provides nonvanishing masses except for $\Pi_3'^\pm$ and $\Pi_3'^0$, which are eaten by gauge bosons $W^\pm$ and Z^0. For H^0 states, the potential (4) provides nonvanishing masses only for the states $H_2'^0$ and $H_3'^0$. We do not consider any gauge bosons which eat the massless scalar fields $H_1'^0$ and H_i^{j0} ($i \neq j$), so that the scalar bosons can appear as physical massless Higgs bosons (hereafter, we denote them as h^0). Since the massless bosons h^0 cannot couple with $\overline{f}f$, $A_\mu A^\mu$, $W_\mu^- W^{+\mu}$ and $Z_\mu Z^\mu$ at tree level, they are harmless to phenomenology in the present accelerator experiments.

The massive Higgs boson $H_3'^0$ corresponds to the physical neutral Higgs boson in the standard model, and it can couple with $W_\mu^- W^{+\mu}$ and $Z_\mu Z^\mu$. We suppose that massive physical bosons $\Pi^\pm$, Π^0, $H_2'^0$ and $H_3'^0$ are heavier than weak bosons $W^\pm$ and Z^0. Therefore, a typical production mode of our physical Higgs bosons is $e^+ + e^- \to Z^* \to H_3'^0 + Z^* \to H_3'^0 + f + \overline{f}$, where $f\overline{f}$ denotes $b\overline{b}$, $\tau\overline{\tau}$, $\mu\overline{\mu}$, and so on, and Z^* means a virtual Z^0. Since we suppose that heavy fermions F are extremely heavy in comparison with weak bosons and massive physical Higgs bosons ϕ_L, our physical Higgs bosons can decay neither into $f\overline{f}$ (at tree level) nor into $f\overline{F}$ ($F\overline{f}$), so that the dominant decay modes of $H_3'^0$ are $H_3'^0 \to h^0\overline{h}^0$ ($h^0 = H_1'^0, H_i^{j0}$). As a result, we will observe a characteristic production mode of $H_3'^0$;

$$e^+ + e^- \to Z^* \to Z^* + H_3'^0 \to f\overline{f} + (\text{neutral particles}) \ .$$

For more phenomenology of the physical Higgs bosons $\Pi^\pm$, Π^0 and H^0, we will discuss elsewhere.

In conclusion, we have proposed a unified mass matrix model of quarks and leptons with seesaw-type matrix form (2) on the basis of a Higgs mechanism scenario for U(3)-family nonet bosons. The Higgs potential (4) can lead to the charged lepton mass relation (1) when we suppose $O_e = \mathbf{1}$, which is provided by the U(3)-family singlet Higgs boson Φ_0. On the other hand, the matrix form O_q in quark mass matrix M_q, (8), has been chosen from a phenomenological consideration. (We will need further plausible explanation on the reason why there is such the Higgs boson Φ_X which couples only with S_3 symmetric states of F_i.) Then, quark mass ratios and KM matrix elements can be fitted only by three parameters b_u, b_d and β_d fairly well. It is worth while that we can obtain a large ratio of m_t/m_b together with a reasonable ratio m_u/m_d without taking so hierarchically different values of b_u and b_d, i.e., with taking $b_d/b_u \simeq 3$. It should also be noted that since our $SU(2)_L$ doublet Higgs bosons ϕ_L couple only between the ordinary fermions f_L and the heavy fermions F_R, the physical Higgs bosons ϕ_L cannot decay into two ordinary fermions at tree level. If our model is true, it will require re-investigation in experimental search for physical Higgs bosons.

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[11] We do not consider a grand unification scenario. We tacitly suppose a composite picture for Higgs bosons and fermions. The scale Λ_X means a scale at which the Higgs bosons Φ_0 and Φ_X have their large VEV's, and the U(3)-family singlet boson Φ_0 couples with all heavy fermions F universally.

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[13] GALLEX collaboration, P. Anselman *et al.*, Phys. Lett. **B285**, 390 (1992).